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R Parity in Standard-like Superstring Models

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ABSTRACT

We investigate the R symmetries of standard-like superstring models. At the level of the cubic superpotential there are three global $U(1)$ R symmetries. These are broken explicitly by $N > 3$ terms in the superpotential and spontaneously by scalar VEVs necessary to preserve supersymmetry at M_P . A Z_2 discrete symmetry remains but is equivalent to fermion number modulo 2. These models possess an effective R parity which arises from the interplay between the gauged $U(1)_{B-L}$ and $U(1)_{r_{j+3}}$.

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1. Introduction

Supersymmetric extensions of the Standard Model (SM), whatever their origin, may contain dangerous terms which violate baryon number (B) and lepton number (L) at unacceptably large rates. This is because, contrary to the SM case, in these models the SM gauge symmetry does not result in accidental (global) B and L symmetries. For example, the superpotential of the minimally supersymmetric Standard Model (MSSM) may contain (dimension four) terms like [1,2]

$$c_1 Q_i L_j d_k + c_2 u_i d_j d_k + c_3 L_i L_j e_k, \quad (1)$$

where i, j, k are generation indices. The first two operators induce an unacceptably large proton decay rate unless $c_1 c_2 < 10^{-26}$ [2] (for squark and slepton masses of at most a TeV). The third operator induces L violating rare events such as $\mu \rightarrow e^+ e^- e^+$ unless c_3 is sufficiently suppressed [3]. It is difficult to explain such small couplings without resorting to a symmetry that protects them. An elegant way of obtaining such small (or vanishing) couplings is to impose an R parity on the model [1,2]. This is a Z_2 symmetry which does not commute with supersymmetry (SUSY) and under which the superfields have the following charges: Q_i, L_i, u_i, d_i, e_i (and N_i , the right-handed neutrino if it exists) all have -1 (odd), H_1, H_2 and the vector superfields have $+1$ (even). In other words, under R parity all SM states are even whereas all their superpartners are odd. R parity, as defined above, eliminates the dangerous terms in Eq. (1) without affecting the usual terms in the MSSM Lagrangian.

The above considerations hold not only for MSSM but for any supersymmetric extension of the SM. In particular, any superstring model which reduces to MSSM or some extension of it at the TeV scale must somehow suppress the terms in Eq. (1) enough so that constraints from B and L violation are satisfied. In this letter, we examine the R parity in standard-like superstring models[4,5]. We find that there are three continuous $U(1)$ R symmetries at the cubic level of the superpotential. These are broken explicitly by higher order ($N > 3$) terms in the superpotential

and spontaneously by the scalar VEVs which are necessary in order to preserve SUSY around the Planck scale. Then, only a discrete Z_2 subgroup of the three $U(1)$'s survive. We find that this Z_2 is not the R parity usually assumed in MSSM and does not eliminate any of the terms in Eq. (1). In fact, it is equivalent to fermion number modulo 2. On the other hand, standard-like superstring models possess an effective R parity which arises from the particular charges of observable and hidden sector states under gauged $U(1)_{B-L}$ and $U(1)_{r_{j+3}}$.

The standard-like superstring models that we consider have the following properties [4,5]:

1. $N = 1$ space-time SUSY.
2. A $SU(3)_C \times SU(2)_L \times U(1)^n$ hidden gauge group.
3. Three generations of chiral fermions and their superpartners, with the correct quantum numbers under $SU(3)_C \times SU(2)_L \times U(1)_Y$.
4. Higgs doublets that can produce realistic electro-weak symmetry breaking.
5. Anomaly cancellation, apart from a single ‘‘anomalous’’ $U(1)$ which is canceled by application of the Dine–Seiberg–Witten (DSW) mechanism [6].

The standard-like superstring models are constructed in the four dimensional free fermionic formulation [7]. The models are generated by a basis of eight boundary condition vectors for all world-sheet fermions [4,5]. The observable and hidden gauge groups after application of the generalized GSO projections are $SU(3)_C \times U(1)_C \times SU(2)_L \times U(1)_L \times U(1)^{6*}$ and $SU(5)_H \times SU(3)_H \times U(1)^2$, respectively. The weak hypercharge is given by $U(1)_Y = \frac{1}{3}U(1)_C + \frac{1}{2}U(1)_L$ and has the standard $SO(10)$ embedding. The orthogonal combination is given by $U(1)_{Z'} = U(1)_C - U(1)_L$.

The models have six right-handed gauge $U(1)_r$ symmetries which correspond to the right-handed world-sheet currents $\bar{\eta}^j \eta^{j*}$ ($j=1,2,3$) and $\bar{y}^3 \bar{y}^6, \bar{y}^1 \bar{\omega}^5, \bar{\omega}^2 \bar{\omega}^4$. At

* $U(1)_C = \frac{3}{2}U(1)_{B-L}$ and $U(1)_L = 2U(1)_{T_{3R}}$.

the level of the cubic superpotential, there is a left-handed global $U(1)_\ell$ symmetry for every right-handed gauge $U(1)_r$. The six left-handed global symmetries correspond to the left-handed world-sheet currents $\chi^{12}, \chi^{34}, \chi^{56}$ and $y^3 y^6, y^1 \omega^5, \omega^2 \omega^4$. We concentrate on the first three of these since their sum gives the $U(1)$ of the $N = 2$ world-sheet SUSY [8] algebra which means that $\chi^{12}, \chi^{34}, \chi^{56}$ are SUSY charges. This can also be seen from the basis vector S

$$S = (\underbrace{1, \dots, 1}_{\psi^\mu, \chi^{12,34,56}}, 0, \dots, 0 | 0, \dots, 0). \quad (2)$$

S plays the part of the SUSY generator in the sense that the SUSY partners of the states from any sector α are given by the sector $S + \alpha$. The three χ^{ij} 's make up the SUSY generator S . Therefore different (scalar or fermionic) components of superfields will have different Q_{ℓ_i} ($i = 1, 2, 3$) charges exactly as for R symmetries. We will see that the world-sheet currents $\chi^{12}, \chi^{34}, \chi^{56}$ correspond to space-time global $U(1)$ R symmetries at the cubic level of the superpotential. Of course, due to a well-known theorem [9], there are no continuous global symmetries in strings. These global $U(1)$ R symmetries are broken explicitly by $N > 3$ terms in the superpotential and spontaneously by scalar VEVs required by the stability of the SUSY vacuum near the Planck mass. As a result there remains only a Z_2 subgroup (R parity of the model) which is equivalent to fermion number modulo 2.

A generic standard-like superstring model including the complete massless spectrum with the gauge quantum numbers and the cubic superpotential were presented in Ref. [4] and will not be repeated here. We use the notation of Ref. [4] throughout this letter. The massless states with their R charges given by $U(1)_{\ell_i}$ are:

(a) The sectors $b_{1,2,3}$ which give the three chiral generations. States from each sector b_i have $Q_{\ell_i} = 1/2$. The superpartners of the chiral fermions in sector b_i come from the sectors $S + b_i$ with $Q_{\ell_j} = Q_{\ell_k} = -1/2$ where $i \neq j \neq k$.

(b) The sector $b_1 + b_2 + \alpha + \beta$ gives the following scalars: a weak doublet h_{45} ,

a color triplet D_{45} and $SO(10)$ singlets $\Phi_{45}, \Phi_1^\pm, \Phi_2^\pm, \Phi_3^\pm$ all with the R charges $Q_{\ell_1} = Q_{\ell_2} = -1/2$. The superpartners of these states have $Q_{\ell_3} = 1/2$.

(c) The Neveu–Schwarz sector gives the graviton, the dilaton, the antisymmetric tensor and the gauge bosons of the model all with vanishing R charges. In addition this sector gives the weak doublets $h_{1,2,3}$ and the singlets $\Phi_{23}, \Phi_{13}, \Phi_{12}$ and ξ_1, ξ_2, ξ_3 with the R charges:

$$Q_{\ell_1}(h_1) = Q_{\ell_1}(\Phi_{23}) = Q_{\ell_1}(\xi_1) = -1, \quad (3a)$$

$$Q_{\ell_2}(h_2) = Q_{\ell_2}(\Phi_{13}) = Q_{\ell_2}(\xi_2) = -1, \quad (3b)$$

$$Q_{\ell_3}(h_3) = Q_{\ell_3}(\Phi_{12}) = Q_{\ell_3}(\xi_3) = -1, \quad (3c)$$

and all other charges vanish. The supersymmetric partners of these are obtained by substituting the degenerate vacuum of the S sector instead of the Neveu–Schwarz vacuum. The superpartners of the graviton, the dilaton, the antisymmetric tensor and the gauge bosons have $Q_{\ell_1} = Q_{\ell_2} = Q_{\ell_3} = 1/2$. The superpartners of other Neveu–Schwarz states with $Q_{\ell_i} = -1$ have $Q_{\ell_j} = Q_{\ell_k} = 1/2$, $Q_{\ell_i} = -1/2$ where again $i \neq j \neq k$.

The barred counterparts of the above states (in the notation of Ref. [4]) have the same $U(1)_{\ell_i}$ charges as the unbarred states. In addition there are hidden sector states from sectors $b_i + 2\gamma$ and $b_{1,2} + b_3 + \alpha \pm \gamma + (I)$ which have nonzero Q_{ℓ_i} [4,5]. None of the results we obtain change if these states are taken into account and therefore we will neglect them in the following.

We see that the states of the model and their superpartners have different charges under the three R symmetries. This is expected since the particles and their SUSY counterparts appear in the same superfields and the superspace parameter θ carries R charges. We find the R charges of θ to be $Q_{\ell_1} = Q_{\ell_2} = Q_{\ell_3} = 1/2$ by inspecting the difference between the R charges of the different components of superfields. The cubic superpotential is obtained by calculating the cubic world-sheet correlators $A_3 \sim \langle V_1^f V_2^f V_3^b \rangle$ using the rules of Ref. [8] and is given in Ref. [4].

(Here V_i^f (V_i^b) are the fermionic (scalar) components of the vertex operators.) All of the terms in the cubic superpotential have the R charges $Q_{\ell_1} = Q_{\ell_2} = Q_{\ell_3} = -1$ so that the integral over superspace, $\int d^2\theta W$, is R invariant. On the other hand, string selection rules impose the space-time (or field theory) selection rules [8]

$$\sum Q_{\ell_1} = \sum Q_{\ell_2} = \sum Q_{\ell_3} = 0, \quad (4)$$

on the F terms obtained at the cubic level. These are seen, at the field theory level, as symmetries of the cubic superpotential. Since there are no gauge bosons corresponding to these symmetries in the massless spectrum, Q_{ℓ_i} are global symmetries. (The corresponding gauge bosons are necessarily projected out by the generalized GSO projection in order to get $N = 1$ space-time SUSY.)

We now show that the three $U(1)$ R symmetries are explicitly and spontaneously broken. The R symmetries are broken explicitly by $N > 3$ nonrenormalizable contributions to the superpotential. These terms are obtained by calculating correlators between vertex operators [8]

$$A_N \sim \langle V_1^f V_2^f V_3^b \cdots V_N^b \rangle. \quad (5)$$

The nonvanishing terms are obtained by applying the rules of Ref. [8]. In order to obtain the correct ghost charge, for an order N term, $N - 3$ vertex operators are picture changed by taking

$$V_{q+1}(z) = \lim_{w \rightarrow z} e^c(w) T_F(w) V_q(z), \quad (6)$$

where T_F is the world-sheet super current given by

$$T_F = \psi^\mu \partial_\mu X + i \sum_{I=1}^6 \chi^I y^I \omega^I = T_F^0 + T_F^{-1} + T_F^{+1}, \quad (7)$$

with

$$T_F^{-1} = e^{-i\chi^{12}} \tau_{12} + e^{-i\chi^{34}} \tau_{34} + e^{-i\chi^{56}} \tau_{56} \quad T_F^{-1} = (T_F^{+1})^*, \quad (8)$$

where $\tau_{ij} = \frac{i}{\sqrt{2}}(y^i \omega^j + i y^j \omega^i)$ and $e^{i\chi^{ij}} = \frac{1}{\sqrt{2}}(\chi^i + i\chi^j)$. It can be shown that only

the T_F^{+1} piece of T_F contributes to A_N [8] and therefore each picture changing adds one unit to either one of one of the R charges. Now, the string selection rules require that $\sum Q_{\ell_i} = 0$ for $i = 1, 2, 3$ are satisfied after all picture changings have been performed. As a result, a generic order N term which requires $N - 3$ picture changings does not satisfy the conservation rules given in Eq. (4) with the R charges given before. Conversely, only order N terms with $\sum Q_{\ell_1} + Q_{\ell_2} + Q_{\ell_3} = 3 - N$ survive the string selection rules. Thus, $N > 3$ terms, in general, break the R symmetries explicitly. For example at order $N = 5$ we have the quark mass terms [10,11]

$$u_1 Q_1 \bar{h}_1 \bar{\Phi}_i^+ \bar{\Phi}_i^-, \quad (9a)$$

$$d_1 Q_1 h_{45} \Phi_1^+ \xi_2, \quad (9b)$$

$$d_2 Q_2 h_{45} \bar{\Phi}_2^- \xi_1, \quad (9c)$$

which break Q_{ℓ_1} and Q_{ℓ_2} explicitly. Q_{ℓ_3} is explicitly broken by the leptoquark-quark mixing term [12]

$$d_3 D_{45} N_3 \Phi_{13} \Phi_3^+ (\xi_1 + \xi_2). \quad (10)$$

As stated above, the massless sector of the model has an anomalous $U(1)$ gauge symmetry, $U(1)_A$. This anomaly is cancelled by a Green-Schwarz counterterm which induces a Fayet-Iliopoulos term into the D constraints for $U(1)_A$ [6]. The set of F and D constraints is given by the following equations:

$$D_A = \sum_k Q_k^A |\chi_k|^2 = \frac{-g^2 e^{\phi_D}}{192\pi^2} \text{Tr}(Q_A) \frac{1}{2\alpha'}, \quad (11a)$$

$$D'^j = \sum_k Q_k'^j |\chi_k|^2 = 0 \quad j = 1 \dots 5, \quad (11b)$$

$$D^j = \sum_k Q_k^j |\chi_k|^2 = 0 \quad j = C, L, 7, 8, \quad (11c)$$

$$W = \frac{\partial W}{\partial \eta_i} = 0, \quad (11d)$$

where χ_k are the fields that get VEVs and Q_k^j are their charges. W is the cubic superpotential and η_i are the fields which do not get VEVs. α' is the string tension and $Tr(Q_A) = 180$ in this model. Eq. (11a) is the D constraint for the anomalous $U(1)_A$. We see that some $SO(10)$ singlet scalars must get Planck scale VEVs in order to satisfy (11a) and preserve SUSY around the Planck scale. Then, due to the other F and D constraints most or all of the other scalars also obtain VEVs. Since these scalars in general have nonzero R charges all three R charges are spontaneously broken around the Planck scale. For example, in the model of Ref. [4] under consideration, Φ_{45} must get a VEV in order to satisfy Eq. (11a). This and other VEVs of scalars coming from the sector $b_1 + b_2 + \alpha + \beta$ break Q_{ℓ_1} and Q_{ℓ_2} spontaneously around the Planck scale. SUSY F constraints in the observable sector require that $\langle \Phi_{12} \rangle = \langle \bar{\Phi}_{12} \rangle = \langle \xi_3 \rangle = 0$ [10]. Thus, Q_{ℓ_3} can only be broken by hidden sector VEVs. For example, $\langle V_1 \rangle$ and $\langle \bar{V}_2 \rangle$ which are needed to get quark mixing break Q_{ℓ_3} spontaneously around the Planck scale [11].

After the three $U(1)$ R symmetries are broken explicitly and spontaneously, there still remains a discrete R symmetry in the model. This is a Z_2 subgroup (of $U(1)^3$), i.e. an R parity. Charges of states under this R parity are given by $\exp(i\pi Q_R)$ where $Q_R = 2 \sum Q_{\ell_i} + Q_{\ell_3}$. The Z_2 charge defined this way is invariant under picture changing since every picture changing operation changes the sum by one and Q_R by two. Q_R is not broken spontaneously either since all $SO(10)$ singlet scalars which get VEVs have $Q_R = 0, 2$. We find that under this Z_2 all superfields are even and θ is odd. As a result, all scalars are even and all fermions are odd. This is equivalent to fermion number modulo 2 and therefore does not give any new constraints on the model. In particular this Z_2 symmetry which is not the usual R parity (under which matter fermions are even) does not eliminate the dangerous terms in Eq. (1).

In light of this result, how can the terms in Eq. (1) be suppressed in standard-like superstring models? In these models $B - L$ is gauged and given by $U(1)_{B-L} = 2U(1)_C/3$. Conservation of $U(1)_{B-L}$ (or equivalently of $U(1)_C$) eliminates the terms in Eq. (1) since they violate $B - L$ in addition to B and L . Order $N > 3$

terms which can induce these operators by scalar VEVs exist but they are all proportional to $\langle \tilde{N}_i \rangle$, the VEVs of the right-handed sneutrinos. Explicitly the nonrenormalizable terms in the superpotential are [10]

$$(u_3 d_3 + Q_3 L_3) d_2 N_2 \Phi_{45} \bar{\Phi}_2^-, \quad (12a)$$

$$(u_3 d_3 + Q_3 L_3) d_1 N_1 \Phi_{45} \Phi_1^+, \quad (12b)$$

$$u_3 d_2 d_2 N_3 \Phi_{45} \bar{\Phi}_2^- + u_3 d_1 d_1 N_3 \Phi_{45} \Phi_1^+, \quad (12c)$$

$$Q_3 L_1 d_3 N_1 \Phi_{45} \Phi_3^+ + Q_3 L_1 d_1 N_3 \Phi_{45} \Phi_3^+, \quad (12d)$$

$$Q_3 L_2 d_3 N_2 \Phi_{45} \bar{\Phi}_3^- + Q_3 L_2 d_2 N_3 \Phi_{45} \bar{\Phi}_3^-, \quad (12e)$$

for the (effective dimension four) B violating operators. L violating operators have the same generic form and will not be written explicitly.

Now, one might think that, in these models, since $B - L$ is gauged, as long as it is not spontaneously broken by $\langle \tilde{N}_i \rangle$ the dangerous terms in Eq. (1) are eliminated. The situation is more complicated since there are hidden sector states with nonzero $B - L$ charge which can get VEVs. In the notation of Ref. [4], these are H_i $i = 14, \dots, 26$ with $Q_{B-L} = \pm 1/2$. SUSY F constraints in the hidden sector require most of the VEVs of H_i to vanish [12]. Still, either H_{23}, H_{25} or H_{24}, H_{26} may get VEVs and break $U(1)_{B-L}$ spontaneously near the Planck scale. In addition, the pair H_{24}, H_{26} has the correct Q_{B-L} charge to render the terms of Eq. (1) neutral under $B - L$. There is no danger of having B (or L) violating terms containing $H_{24} H_{26}$ however, due to conservation of local $U(1)_{r_{j+3}}$. (Strictly speaking, this is true only when H_{19} and H_{20} get Planck scale masses and decouple. In these models, they get masses of 10^{17} GeV from the cubic superpotential [4].) Therefore, even though gauged $B - L$ can be broken by hidden sector VEVs, the terms in Eq. (12a-e) are still the only ones to induce those in Eq. (1). This is due to a conspiracy between the $U(1)_{B-L}$ and $U(1)_{r_{j+3}}$ charges of the massless states in the observable and hidden sectors. One can think of this as an effective R parity of the model which can only be broken by nonvanishing $\langle \tilde{N}_i \rangle$. It is this effective R parity which requires the states N_i to appear in all the B (and L)

violating terms in Eqs. (12a-e). The magnitude of the coefficients c_i in Eq. (1) is controlled by $\langle \tilde{N}_i \rangle$. Since there are no other constraints on $\langle \tilde{N}_i \rangle$ (from SUSY etc.), one can choose them to be zero or small enough to satisfy the constraints from proton lifetime (and L violating processes). Reversing the argument, proton lifetime gives an upper bound on $\langle \tilde{N}_i \rangle$. The strongest constraints arise from the terms in Eq. (12a,b) and give $\langle \tilde{N}_i \rangle / M < 10^{-11}$ or $\langle \tilde{N}_i \rangle < 10^7 \text{ GeV}$ for TeV scale squarks.

To summarize, we find that there are three continuous $U(1)$ R symmetries at the cubic level of the superpotential. These are broken explicitly by higher order ($N > 3$) terms in the superpotential and spontaneously by the scalar VEVs which are necessary in order to preserve SUSY around the Planck scale. Then, only an R parity (Z_2) which is equivalent to fermion number modulo 2 survives. On the other hand, standard-like superstring models possess an effective R parity which arises from the particular charges of observable and hidden sector under gauged $U(1)_{B-L}$ and $U(1)_{r_{j+3}}$.

The effective R parity which can be broken only by $\langle \tilde{N}_i \rangle$ is not enough to rule out large B (or L) violation in standard-like superstring models. There may be effective B (or L) violating terms other than the ones in Eq. (1). For example one must make sure that the effective $N = 4$ terms in the superpotential such as

$$c_4 Q_i Q_j Q_k L_l + c_5 u_i u_j d_k e_l + c_6 Q_i Q_j Q_k H_l \quad (13)$$

which induce dimension five B violating terms are sufficiently suppressed. (Here i, j, k, l are generation indices.) In addition, intermediate scale color triplets (such as the leptoquark D_{45} in standard-like superstring models [12]) can induce effective dimension five or six B violating operators unless they have weak enough couplings to matter fermions and/or large enough masses. A more detailed study shows that B violation arising from the sources mentioned above can be suppressed sufficiently in standard-like superstring models so that constraints from proton lifetime are satisfied [14].

Arguments similar to the ones used in this letter can also be applied to the other global symmetries of the cubic superpotential. For example, in the model under consideration, the left-handed world-sheet currents $y^3 y^6, y^1 \omega^5, \omega^2 \omega^4$ give three other global $U(1)$ symmetries of the cubic superpotential. These too are broken explicitly by $N > 3$ terms and spontaneously by the scalar VEVs and may lead to discrete symmetries. In addition, string selection rules together with the specific charges of massless states may result in discrete symmetries which are not subgroups of continuous global symmetries.

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REFERENCES

1. S. Dimopoulos and H. Georgi, Nucl. Phys. B **193** (1981) 150.
2. S. Weinberg, Phys. Rev. D **26** (1982) 287; N. Sakai and Y. Yanagida, Nucl. Phys. B **197** (1982) 533.
3. L. J. Hall and M. Suzuki, Nucl. Phys. B **231** (1984) 419.
4. A.E. Faraggi, Phys. Lett. B **278** (1992) 131.
5. A.E. Faraggi, Phys. Lett. B **274** (1992) 47; Nucl. Phys. B **387** (1992) 289.
6. M. Dine, N. Seiberg and E. Witten, Nucl. Phys. B **289** (1987) 585; J.J. Atick, L.J. Dixon and A. Sen, Nucl. Phys. B **292** (1987) 109; S. Cecotti, S. Ferrara and M. Villasante, Int. J. Mod. Phys. A **2** (1987) 1839.
7. I. Antoniadis, C. Bachas, and C. Kounnas, Nucl. Phys. B **289** (1987) 87; I. Antoniadis and C. Bachas, Nucl. Phys. B **298** (1988) 586; H. Kawai, D.C. Lewellen, and S.H.-H. Tye, Phys. Rev. Lett. **57** (1986) 1832; Phys. Rev. D **34** (1986) 3794; Nucl. Phys. B **288** (1987) 1.

8. S. Kalara, J. Lopez and D.V. Nanopoulos, Phys. Lett. B **245** (1991) 421;
Nucl. Phys. B **353** (1991) 650.
9. T. Banks and L. Dixon, Nucl. Phys. B **307** (1988) 93.
10. A.E. Faraggi, Nucl. Phys. B **403** (1993) 101.
11. A. E. Faraggi and E. Halyo, Phys. Lett. B **307** (1993) 305; WIS-93/34/MAR-
PH. to appear in Nucl. Phys. B.
12. E. Halyo, WIS-93/114/DEC-PH.
13. A.E. Faraggi, Nucl. Phys. B **407** (1993) 57.
14. E. Halyo, WIS preprint in preparation.